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Warping the Weak Gravity Conjecture

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ABSTRACT

The Weak Gravity Conjecture, if valid, rules out simple models of Natural Inflation by restricting their axion decay constant to be sub-Planckian. We revisit stringy attempts to realise Natural Inflation, with a single open string axionic inflaton from a probe D-brane in a warped throat. We show that warped geometries can allow the requisite super-Planckian axion decay constant to be achieved, within the supergravity approximation and consistently with the Weak Gravity Conjecture. Preliminary estimates of the brane backreaction suggest that the probe approximation may be under control. However, there is a tension between large axion decay constant and high string scale, where the requisite high string scale is difficult to achieve in *all* attempts to realise large field inflation using perturbative string theory. We comment on the Generalized Weak Gravity Conjecture in the light of our results.

1. Introduction

Cosmological inflation stands strong as the leading mechanism to provide the seeds that gave rise to the large structure we observe today in the Universe. Precision observations in the Cosmic Microwave Background provide a window into this very early history of the Universe. The latest results from Planck [1] are in perfect agreement with the simplest inflationary models, driven by the dynamics of a single scalar field rolling down a very flat potential. Current bounds from Planck/BICEP2 on the scalar to tensor ratio in the CMB power spectrum are $r \lesssim 0.12$ (95% CL). Any future detection of tensor modes would have the remarkable implications, via the Lyth relation [2–4], that inflation occurred at scales close to the Planck scale:

$$V_{inf}^{1/4} \approx \left(\frac{r}{0.1}\right)^{1/4} \times 1.8 \times 10^{16}\,\text{GeV} \tag{1}$$

and that the inflaton field had super-Planckian excursions:

$$\frac{\Delta\phi}{M_{Pl}} \gtrsim 0.25 \times \left(\frac{r}{0.01}\right)^{1/2}. \tag{2}$$

"Large field" inflationary models are intriguing not only due to their robust prediction of high scale inflation with observable primordial gravitational waves. They depend sensitively on the degrees of freedom comprising the ultraviolet completion of gravity. In particular, Planck suppressed corrections to the slow-roll inflaton potential typically become large when the inflaton varies over super-Planckian scales. One idea to protect the slow-roll inflaton potential from dangerous quantum corrections is to invoke a shift symmetry in the inflaton field, for instance by identifying the inflaton with a Goldstone boson, the axion. The classical example in this vein is Natural Inflation [5], now tightly constrained by the latest CMB observations.[1] In Natural Inflation, the axion enjoys a continuous shift symmetry within the perturbative approximation. This is broken to a discrete symmetry by non-perturbative effects, which generate a potential of the form:

$$V(\phi) = V_0\left(1 \pm \cos\left(\frac{\phi}{f}\right)\right), \tag{3}$$

where f is the axion decay constant. The potential is sufficiently flat for slow-roll inflation provided that $f \gg M_{Pl}$, and as a consequence the axion can undergo super-Planckian field excursions.

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[1] Although massive modes during inflation can change the classic NI predictions by generating a smaller than unity speed of sound bringing the model back to the allowed parameter region, as shown in [6].

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The current bound on the Natural Inflation potential (3), given by PLANCK from the spectral index, n_s, is $f/M_{Pl} \gtrsim 6.8$ (95% CL) [1].

To understand whether or not such ideas are viable requires the embedding of large field models in a theory of quantum gravity. In fact, much interest has recently been generated by the possibility that general features of quantum gravity can constrain inflationary models with observable consequences. The Weak Gravity Conjecture [7] roughly proposes that gravity must be the "weakest force" in a quantum theory of gravity, in order to avoid stable black hole remnants. For example, for a four dimensional theory describing gravity and a $U(1)$ gauge sector with gauge coupling, $\mathbf{g}$, there must exist a state with mass, m, which satisfies $m \lesssim M_{Pl}\,\mathbf{g}$. Moreover, the effective field theory has a new UV cutoff scale, $\Lambda \sim M_{Pl}\,\mathbf{g}$. This conjecture was used [7] to rule out Extra Natural Inflation [8], where the inflaton arises from a Wilson line in a five-dimensional $U(1)$ gauge theory compactified on a circle.

The Weak Gravity Conjecture might also be generalized to D dimensions, p-form Abelian gauge fields, and their p spacetime dimensional charged objects. Then, in a four dimensional gravitational theory with a 0-form axion, there must exist an instanton with action, $S_{cl} \lesssim M_{Pl}/f$, where f is the axion decay constant. Although this conjecture lacks convincing motivation from black hole physics, the same phenomenon was observed in [9] in several diverse string theoretic setups. If valid, it would essentially rule out single field models of inflation with super-Planckian axion decay constants, as instanton corrections would always introduce higher harmonics to the inflation potential:

$$e^{n\frac{M_{Pl}}{f}} e^{in\theta} , \tag{4}$$

effectively limiting the axion field range. Analogous arguments considering gravitational instantons in effective field theory lead to similar conclusions [10]. Moreover, the examples studied in [9] demonstrated the difficulty in obtaining axions with large decay constants within the limits of perturbative string theory, and raised the question if this is possible at all.

Recent work has focused on whether these constraints from the Generalized Weak Gravity Conjecture on axion inflation can be evaded, in particular, by introducing multiple axion fields [11–19]. There has also been a large amount of work towards developing string theoretic models of large field inflation with sub-Planckian axion decay constants. These come under two main classes, firstly, stringy mononomial chaotic inflation scenarios where monodromy effects explicitly break the axion shift symmetry [20–28], and secondly, many field models where multiple axions generate an effective decay constant that is super-Planckian [29–34]. Most constructions have used closed string axions in type II string theory.

In this letter, we revisit open string, single field models of axion inflation, in the light of quantum gravity constraints and the Weak Gravity Conjecture discussed above. Open string inflatons include Wilson lines on wrapped Dp-branes and the position moduli of moving Dp-branes. In fact, these scenarios are T-dual to each other. By considering D3-branes moving down a long warped throat, Baumann and McAllister pointed out a rigid, sub-Planckian upper bound on the field range [35,36], and this can indeed be interpreted as a consequence of the Weak Gravity Conjecture. However, single field models have been proposed with moderately super-Planckian decay constants. Wilson lines on wrapped Dp-branes with sub-Planckian decay constants were studied in [37], and with super-Planckian decay constants in [38]. Planckian decay constants from wrapped Dp-branes moving down or around a warped throat were found, respectively, in [39] and [40].

We show that warping and wrapped volumes indeed allow for single field models with super-Planckian axion decay constant consistently with the Weak Gravity Conjecture. The large decay constants are generated within the perturbative limits of the supergravity approximation, and initial estimates support the validity of the probe brane approximation used. Moreover, scalar potentials that break the continuous axion shift symmetry to a discrete one are potentially generated by classical or loop effects [8,40–43]. Therefore, any non-perturbative instanton effects are by construction exponentially suppressed, and would not rule out single field, large field, slow-roll inflation.

Unfortunately, these results do not lead to promising models of large field inflation and observable primordial gravitational waves. This is because in explicit constructions, there is a tension between obtaining large decay constant and high string scale. In fact, as we emphasize, it is always difficult to obtain a sufficiently high string scale within the limits of perturbation theory. This presents an important challenge in building string theoretic models of large field inflation.

The paper is organized as follows. In the next section we introduce inflation from D-branes in warped geometries and fix our conventions. In Section 3 we study scenarios in which the candidate axionic inflaton is a Wilson line on a wrapped D-brane, and in Section 4 we turn to the T-dual picture of the position modulus of a wrapped D-brane. The most attractive scenario studied can be found at the end of this section. Finally, in Section 5 we discuss our results, and the light they shed on the Generalized Weak Gravity Conjecture.

2. Open string inflation

Our starting point is a generic type IIB string warped compactification from ten to four dimensions with metric (in the Einstein frame[2]):

$$ds^2 = h^{-1/2}(r) g_{\mu\nu} dx^\mu dx^\nu + h^{1/2}(r) g_{mn} dx^m dx^n , \tag{5}$$

where $\mu, \nu = 0, \dots 3$, $m, n = 4, \dots 9$ and $h(r)$ is the warp factor, possibly trivial, depending on a radial-like direction in the internal space. For example, for an $adS_5 \times X_5$ geometry which describes well a generic warped throat generated by branes and fluxes in the mid-throat region, we have:

$$h(r) = L^4/r^4 \quad \text{and} \quad ds_6^2 = g_{mn} dx^m dx^n = dr^2 + r^2 d\Omega_5^2 , \tag{6}$$

where L is the adS length scale and $d\Omega_5^2 = \tilde{g}_{ij} d\phi^i d\phi^j$ is the metric on some five-dimensional Einstein–Sasaki space that ensures $\mathcal{N} = 1$ supersymmetry. To construct a smooth, compact internal space out of the adS throat, we take $r_{IR} < r < r_{UV}$, where the adS region is glued to the tip of the throat at the IR cutoff and to a compact Calabi–Yau at the UV cutoff [44].

The four-dimensional Planck mass after compactification takes the form:

$$M_{Pl}^2 = \frac{4\pi \mathcal{V}_6^w}{g_s^2} M_s^2 \quad \text{with} \quad \mathcal{V}_6^w l_s^6 = \int d^6 y \sqrt{\det g_{mn}}\, h , \tag{7}$$

where we defined the string scale as $l_s^2 = M_s^{-2} = (2\pi)^2 \alpha'$, $g_s = e^{\varphi_0}$ is the string coupling and $\mathcal{V}_6^w$ is the warped volume of the six-dimensional internal space in string units. For example, assuming most of the volume comes from the middle region of an adS throat generated by N D3-branes at its tip, we have:

$$L^4 = \frac{g_s N}{4\mathcal{V}_5} l_s^4 \tag{8}$$

[2] Our conventions for going from string to Einstein frame are $G^E_{MN} = e^{\frac{\varphi_0 - \varphi}{2}} G^s_{MN}$, where φ is the dilaton, whose vev $\langle\varphi\rangle = \varphi_0$ defines the string coupling as $g_s = e^{\varphi_0}$. In these conventions the volumes evaluated in the background are frame independent.

and

$$\mathcal{V}_6^w l_s^6 = \frac{1}{2}\mathcal{V}_5 L^4 r_{UV}^2 \,, \tag{9}$$

with $\mathcal{V}_5 = \int d\Omega_5$ the dimensionless volume of the base of the cone.

To the above space we add a probe space-filling Dp-brane wrapping a $(p-3)$-cycle in the internal space, with worldvolume coordinates ξ^A $(A, B = 0, \ldots, p+1)$ and DBI action (in the Einstein frame):

$$S_{DBI} = -T_p \int d^{p+1}\xi\, \sqrt{-\det(\gamma_{AB} + \mathcal{F}_{AB})}\,. \tag{10}$$

The tension of the brane in the Einstein frame is given by:

$$T_p = \mu_p\, g_s^{-1} \quad \text{with} \quad \mu_p = (2\pi)^{-p}(\alpha')^{-\frac{(p+1)}{2}}\,. \tag{11}$$

Also, $\gamma_{AB} = g_{MN}\partial_A X^M \partial_B X^N$ is the pullback of the ten-dimensional metric onto the brane $(M, N = 0, \ldots 9)$. Finally, $\mathcal{F}_{AB} = \mathcal{B}_{AB} + 2\pi\alpha' F_{AB}$, with $\mathcal{B}_{AB}$ the pullback of the NSNS 2-form onto the brane, and F_{AB} the field strength associated to the worldvolume gauge field. We will choose static coordinates for the brane, so $\xi^A = (x^\mu, y^a)$ with y^a the $(p-3)$ internal coordinates along the brane. We will comment on the validity of the probe approximation below.

Possible open string inflaton fields in the above system include Wilson line moduli associated with the worldvolume gauge field [37,38] and the moduli describing the position of the Dp-brane in the compact space [45–47]. We will consider single field, large field inflation and the Weak Gravity Conjecture in these setups, which are related to each other by T-duality.

3. Wilson line inflation

When the $(p-3)$-cycle wrapped by the Dp-brane contains a non-trivial 1-cycle,[3] parameterized by some coordinate ϕ, the brane can have a Wilson line wrapping its worldvolume:

$$e^{i\theta} = e^{i\oint A_\phi d\phi}\,. \tag{12}$$

Upon dimensional reduction, the DBI action (10) for the brane in the background (5)–(6) includes a 4D gauge kinetic term for the $U(1)$ worldvolume gauge field:

$$S = -\int d^4x\sqrt{-g}\,\frac{1}{4\mathbf{g}_4^2}\,F_{\mu\nu}F^{\mu\nu}\,, \tag{13}$$

where the four-dimensional effective gauge coupling constant, $\mathbf{g}_4$, is computed to be:

$$\mathbf{g}_4^2 = (2\pi) g_s h_0^{(3-p)/4}\,(n\,\mathcal{V}_{p-3})^{-1}\,. \tag{14}$$

Here, h_0 is the value of the warp factor evaluated at the brane position and $\mathcal{V}_{p-3}$ is the unwarped volume of the $(p-3)$-cycle wrapped by the brane in units of l_s, $\mathcal{V}_{p-3} l_s^{p-3} = \int d^{p-3}y\sqrt{\det g_{ab}}$ (notice that $\mathcal{V}_{p-3}$ also depends on the position of the brane, see eq. (6)). Finally n is the wrapping number.

At the same time, the DBI action leads to a kinetic term for the Wilson line modulus, $\theta = 2\pi A_\phi$, which takes the form

$$S = -\int d^4x\sqrt{-g}\,\frac{f^2}{2}\partial_\mu\theta\partial^\mu\theta \tag{15}$$

where the axion decay constant, f, is computed to be:

$$f^2 = \frac{n\, g_s^{-1}}{(2\pi)^3}\frac{h_0^{(p-7)/4}\mathcal{V}_{p-3}}{R_0^2} \tag{16}$$

Here, $2\pi R_0$ is the unwarped length of the 1-cycle with line element $ds^2 = R_0^2\, d\phi^2 = r_0^2 \tilde{g}_{\phi\phi} d\phi^2$.

The axion, θ, has a continuous shift symmetry descending from the higher dimensional gauge symmetry, which may be broken to a discrete one. A slow-roll potential can thus be generated for the Wilson line axion, θ, by fluxes, brane backreaction, warping, loop contributions and/or other effects, including non-perturbative ones [8,49,37,50]. Large field excursions for θ are encoded in the field's decay constant by $f/M_{Pl} > 1$, with:

$$\frac{f^2}{M_{Pl}^2} = \frac{g_s}{2(2\pi)^4}\frac{h_0^{(p-7)/4} l_s^2}{R_0^2}\frac{n\,\mathcal{V}_{p-3}}{\mathcal{V}_6^w}\,. \tag{17}$$

Notice that for $r_0 > L$, the warp factor can enhance f/M_{Pl}, although the final behaviour of the latter also depends on the wrapped volumes.

As a first example, consider a Dp-brane in an unwarped, possibly anisotropic toroidal orientifold compactification, with j directions of size R, and $6-j$ directions of size L. Taking the brane to wrap a $(p-3)$-cycle with j_b directions of size R (so $j \geq j_b$) and $p-3-j_b$ directions of size L (so $6-j \geq p-3-j_b$), and the Wilson line wrapping one of the R-directions (so $j_b \geq 1$), we have:

$$\frac{f^2}{M_{Pl}^2} = \frac{n\, g_s}{2(2\pi)^4(2\pi)^{9-p}}\left(\frac{l_s}{R}\right)^{2+j-j_b}\left(\frac{l_s}{L}\right)^{9-p-j+j_b}\,. \tag{18}$$

A large f/M_{Pl} would require a substring scale cycle[4] $R < l_s$ and/or $L < l_s$. In this limit the perturbative description breaks down, and the T-dual description should be used.

Next, consider a wrapped D5-brane in an $adS_5 \times X_5$ warped throat. The axion decay constant is then:

$$\frac{f^2}{M_{Pl}^2} = \frac{g_s}{(2\pi)^4}\frac{n\mathcal{V}_2}{\mathcal{V}_5}\frac{l_s^6}{L^6}\frac{r_0^2}{r_{UV}^2}\frac{l_s^2}{R_0^2}\,. \tag{19}$$

Since $\mathcal{V}_2$ goes as r_0^2/l_s^2, the possibility now arises that $f/M_{Pl} > 1$ can be achieved with a brane at the top of a long throat, $r_0 \sim r_{UV}$ and $r_{UV} \gg L$.

Now let us check whether the conditions for large axion decay constant can be fulfilled consistently with the Weak Gravity Conjecture. Indeed, probe D-branes can be used to test the consistency of a string theory configuration by asking whether their worldvolume gauge theories and charged matter satisfy the Weak Gravity Conjecture, wherever they are placed. The Weak Gravity Conjecture implies a new UV cutoff, $\Lambda = M_{Pl}\mathbf{g}_4$, which in the present case is:

$$\Lambda^2 = \frac{(2\pi)^2}{g_s l_s^2}\frac{2}{h_0^{(p-3)/4}}\frac{\mathcal{V}_6^w}{n\,\mathcal{V}_{p-3}}\,. \tag{20}$$

Matter fields charged under the $U(1)$ arise from strings stretching between the Wilson line Dp-brane and a separated parallel Dp-brane. Largest masses arise when the branes are very distant, with $m \sim M_s^2 \ell_\Phi$ and ℓ_Φ the brane separation. So the Weak Gravity Conjecture, $m^2 < M_{Pl}^2\mathbf{g}_4^2$ requires:

$$1 < \frac{(2\pi)^2}{g_s}\frac{l_s^2}{\ell_\Phi^2}\frac{2}{h_0^{(p-3)/4}}\frac{\mathcal{V}_6^w}{n\,\mathcal{V}_{p-3}}\,. \tag{21}$$

[3] Note that this does not require the full six-dimensional internal space to have a 1-cycle, see [48] for examples.

[4] Vanishing cycles do occur in more general string geometries, for example when blowing-up a singularity. The vanishing cycles in the conifold were used in [51] to achieve small axion decay constants from string theory, and in [52] to achieve string axion N-flation models with near Planck scale axion decay constants.

The Weak Gravity Conjecture (21) then imposes an upper limit on f/M_{Pl} (see eq. (17)):

$$\frac{f^2}{M_{Pl}^2} < \frac{h_0^{-1}}{(2\pi)^2}\frac{l_s^4}{\ell_\Phi^2 R_0^2}\,. \tag{22}$$

Considering an $adS_5 \times X_5$ throat (6) as a prototype, the warped length of a string stretching from a D5-brane at r_0 near the top of the throat all the way down to the bottom of the throat is:

$$\ell_\Phi \sim \int_{r_{IR}}^{r_0} h^{1/4} dr = L \ln\left(\frac{r_0}{r_{IR}}\right)\,. \tag{23}$$

Taking $r_{IR} \sim l_s$ and $r_0 \sim r_{UV}$ we then obtain from eq. (21):

$$1 < \frac{(2\pi)^2}{n g_s}\frac{r_0^2/l_s^2}{(\ln(r_0/l_s))^2}\frac{r_0^2\,\mathcal{V}_5}{l_s^2\,\mathcal{V}_2}\,, \tag{24}$$

where again recall that $\mathcal{V}_2$ goes as r_0^2/l_s^2. Note in particular that the Weak Gravity Conjecture does not put an upper limit on r_0, so both the Weak Gravity Conjecture and large axion decay constant could indeed be satisfied with a brane at the top of a long throat.

In the above analysis, we have neglected the backreaction of the wrapped D5-brane onto the background geometry. In fact, a D5-brane would alter the warp factor and internal geometry (cf. (6)), and introduce a non-trivial profile for the dilaton. However, if its contribution to the Einstein equations is much smaller than that of the stack of N D3-branes sourcing the warped throat, then we can safely consider the D5-brane as a probe. The local contribution from a Dp-brane to the traced Einstein's equation goes as [44,39]:

$$\left(T_m^m - T_\mu^\mu\right)^{loc} = (7-p)T_p\Delta^{(9-p)}(\Sigma_{p-3}) \tag{25}$$

where $\Delta^{(9-p)}(\Sigma_{p-3}) = \delta^{(9-p)}(\Sigma_{p-3})/\sqrt{\det g_{9-p}}$ is the covariant delta function on the wrapped $(p-3)$-cycle, Σ_{p-3}, with transverse volume $\int\sqrt{\det g_{9-p}}$. The condition that the backreaction of the wrapped D5-brane is negligible can then be written as:

$$\frac{n}{2N}\frac{T_5}{T_3}\frac{\Delta^{(4)}(\Sigma_2)}{\Delta^{(6)}(\Sigma_0)} \ll 1\,. \tag{26}$$

This condition ensures that the probe approximation is good when the probe is close to the N source D3-branes. We should exercise some caution as the probe becomes very distant from the source branes, at the top of the throat.

For example, for the D5-brane in the $adS_5 \times X_5$ background, we can take the parameters $L \sim \sqrt{3}\,l_s$, $g_s \sim 0.3$, $r_0 \sim r_{UV} \sim 500\,l_s$, $R_0 = r_0\,\tilde{g}_{\phi\phi}^{1/2} \sim l_s$, $n \sim 100$, $\mathcal{V}_2 l_s^2 \sim \pi r_0^2$ and $\mathcal{V}_5 \sim \pi^3$. To compute the backreaction, we estimate transverse volume elements as $\sqrt{\det g_4} \sim h_0 r_0^3 \sin\theta_1$ and $\sqrt{\det g_6} \sim h_0^{3/2} r_0^5 \sin\theta_1 \sin\theta_2$, for some internal coordinate angles θ_1, θ_2. Thus, we can consistently achieve $f \sim 4.2\,M_{Pl}$, $M_s \sim 1.4 \times 10^{-5} M_{Pl}$ and $\Lambda \sim 4.5 \times 10^{-2} M_{Pl}$. Note that we are safely within the limit of perturbative supergravity, but it would be important to study in detail if the probe approximation remains good when the D5-brane is at the top of the long throat. In this sense, the scenario presented at the end of Section 4 will be more reliable. As we will also see there, it is also conceivable that more complicated warped geometries than the adS mid-throat geometry may allow one to achieve even higher f/M_{Pl}.

To summarise, it is plausible that a Wilson line on a wrapped D5-brane can give rise to an axion with large axion decay constant consistently with the Weak Gravity Conjecture. This can be achieved by ensuring a warped throat geometry with a 1-cycle inside the 2-cycle that is wrapped by the D5-brane in a region with small warp factor, $h_0^{1/2} \ll 1$, and thus requires a highly anisotropic long warped throat. A concrete realisation of the proposal would require a throat geometry with these properties. It would likely not, however, correspond to a stringy realisation of the Extra-Natural Inflation proposal [8], as the string scale would probably turn out to be too low. Indeed, note that taking $r_0 \sim r_{UV}$ large to achieve large f/M_{Pl} drives the string scale to smaller values. We will comment further on this below. First we turn to the T-dual description of a D-brane Wilson line, which is a D-brane position field. In this case, more explicit, reliable constructions are possible.

4. Brane position inflation

Once again consider a space-filling Dp-brane wrapping a $(p-3)$-cycle, and now assume that the position of the brane in one of the compact angular dimensions corresponds to an inflaton field. Such scenarios have been studied, for example, in [53,54,39,40,25]. Dimensional reduction of the DBI action (10) gives, in exactly the same way as for the Wilson line scenario, a $U(1)$ gauge group in the four-dimensional effective field theory, see (13), (14). It also leads to a kinetic term for the angular position modulus, θ, of the form (15). The axion decay constant, f, is now given by:

$$f^2 = \frac{(2\pi)\,g_s^{-1}}{l_s^4} h_0^{(p-3)/4} g_{\theta\theta}\, n\, \mathcal{V}_{p-3}\,. \tag{27}$$

A large axion decay constant requires $f/M_{Pl} > 1$ with:

$$\frac{f^2}{M_{Pl}^2} = \frac{1}{2}\, g_s h_0^{(p-3)/4} g_{\theta\theta}\frac{n\,\mathcal{V}_{p-3}}{l_s^2\mathcal{V}_6^w}\,. \tag{28}$$

Now, taking the brane at the tip of the throat, $r_0 \ll L$, we see that the warp factor enhances f/M_{Pl}, although again the wrapped volumes will also affect the overall behaviour.

The axion, θ, corresponds to a periodic direction and so inherits a discrete shift symmetry. Again, the inflaton potential may be induced by fluxes, brane backreaction, warping, loop contributions and/or other effects, including subleading non-perturbative effects. For instance, [40] found a potential of the Natural Inflation form (3), by considering the forces experienced by a D-brane moving in the geometry of the warped resolved conifold glued to a compact Calabi–Yau.

The new UV cutoff implied by the Weak Gravity Conjecture is the same as for the Wilson line scenario, see (20). Charged matter may again arise due to strings stretching to parallel D-branes that are separated by a distance ℓ_Φ along the throat or around the θ-direction. The Weak Gravity Conjecture, $m < \mathbf{g}_4 M_{Pl}$ constraint is the same as before, eq. (21), and implies the following upper bound on the axion decay constant,[5] (28):

$$\frac{f^2}{M_{Pl}^2} < (2\pi)^2\frac{g_{\theta\theta}}{\ell_\Phi^2}\,. \tag{29}$$

We can now consider these constraints in some explicit constructions. Take first a possibly anisotropic unwarped toroidal orientifold compactification, with j directions of large size R and $6-j$ directions of small size L. Consider a Dp-brane wrapping j_b large directions and $p-3-j_b$ small directions, cycling a large direction (so $j_b \le j-1$). The position modulus has axion decay constant given by:

$$\frac{f^2}{M_{Pl}^2} = \frac{n\,g_s}{2}\frac{1}{(2\pi)^{9-p}}\left(\frac{R}{L}\right)^{2+j_b-j}\left(\frac{l_s}{L}\right)^{7-p}\,. \tag{30}$$

[5] We thank Zac Kenton and Steve Thomas for pointing out an error in the previous version of this equation.

So a large decay constant could be achieved with a large angular direction $R \gg L$ for $j - j_b = 1$. We can test the consistency of the required parameters by using the Weak Gravity Conjecture. To this end we introduce another probe Dp-brane, separated by a distance $\ell_\Phi \sim \pi R$ around a large angular direction. The Weak Gravity Conjecture (21) then requires:

$$1 < \frac{8}{n\, g_s}(2\pi)^{9-p}\left(\frac{L}{R}\right)^{2+j_b-j}\left(\frac{L}{l_s}\right)^{7-p} \tag{31}$$

and so imposes an upper bound on the decay constant:

$$\frac{f^2}{M_{Pl}^2} < 4\,, \tag{32}$$

which is at most $f \sim 2\,M_{Pl}$. For example, taking a D3-brane, with $g_s \sim 0.3$, $L \sim l_s$, $R \sim 1.5 \times 10^6\, l_s$, $j = 1$, $j_b = 0$ and $n = 1$, we obtain $f \sim 1.9\, M_{Pl}$, $M_s \sim 2.8 \times 10^{-7} M_{Pl}$ and $\Lambda = 1.4 \times M_{Pl}$. Note that to achieve $f \sim 2\, M_{Pl}$, the string scale is very low.

Let us next consider a Dp-brane moving in a warped throat. The prototypical example is the Klebanov–Strassler warped throat, produced by placing N D3-branes at the tip of the deformed conifold. Away from the conical deformation, in the mid-throat region, the base of the conifold can be taken to be $T^{1,1}$, and the 10D metric takes the form (5), with $h = L^4/r^4$ and

$$\begin{aligned} ds_6^2 = dr^2 &+ \frac{1}{6}r^2\left(d\theta_1^2 + \sin^2\theta_1 d\phi_1^2\right) \\ &+ \frac{1}{6}r^2\left(d\theta_2^2 + \sin^2\theta_2 d\phi_2^2\right) \\ &+ \frac{1}{9}r^2\left(d\psi^2 + \cos\theta_1 d\phi_1 + \cos\theta_2 d\phi_2\right)^2 . \end{aligned} \tag{33}$$

Allowing for a $\mathbb{Z}_k$ orbifolding in the ψ direction, the volume of $T^{1,1}$ is given by $\mathcal{V}_5 = \frac{16}{27k}\pi^3$ and the adS radius of curvature (8) is:

$$L^4 = \frac{27k}{64\pi^3} g_s N l_s^4 . \tag{34}$$

Now consider a D5-brane wrapping a 2-cycle in the throat [6] [39]:

$$r = r_0,\ \psi = \psi_0,\ \theta_1 = -\theta_2,\ \phi_1 = -\phi_2 \tag{35}$$

and moving in an angular direction, say, θ_2.

The condition that the backreaction of the probe D5-brane can be consistently neglected, eq. (26), evaluates in the present case to[7][8]:

$$\frac{n}{12N}\frac{L^2}{l_s^2}\sin\theta_1 \ll 1 . \tag{36}$$

The parameters are also constrained by the Weak Gravity Conjecture. To see how, we introduce another probe Dp-brane, firstly separated from r_0 up the throat to r_{UV}, and secondly separated around the throat along θ_2 at $r = r_0$. We estimate the mass of the corresponding charged matter, $m = M_s^2 \ell_\Phi$, as (cf. (23)):

$$m = M_s^2 L \log\frac{r_{UV}}{r_0}\,, \tag{37}$$

$$m = M_s^2 \frac{\pi L}{\sqrt{6}}\,, \tag{38}$$

where we have taken into account that the lengths of the stretched strings, ℓ_Φ, are warped. The Weak Gravity Conjecture (21) thus implies that (use eqs. (9), (35), $\mathcal{V}_2 l_s^2 = \frac{4\pi}{3} r_0^2$ and $h_0 = L^4/r_0^4$):

$$1 < \frac{32\pi^2}{3}\frac{1}{nk\, g_s}\frac{r_{UV}^2}{l_s^2}\,, \tag{39}$$

where notice that this ensures that the wrapping number, n, and orbifold number, k, are not too large. The Weak Gravity Conjecture thus gives the following upper bound on f:

$$\frac{f^2}{M_{Pl}^2} < \frac{4r_0^2}{L^2}\,. \tag{40}$$

Comparing this with result for the torus (32), we see that warping has relaxed the upper bound on the axion decay constant, and larger super-Planckian values may now be possible within the supergravity approximation.

In fact, one can immediately infer from Eqs. (28), (34), (9), $g_{\theta\theta} = r_0^2/6$, $\mathcal{V}_2 l_s^2 = \frac{4\pi}{3} r_0^2$ and $h_0 = L^4/r_0^4$, that the axion decay constant is given by:

$$\frac{f^2}{M_{Pl}^2} = \frac{3\, nk\, g_s}{8\pi^2}\frac{l_s^2 r_0^2}{L^2 r_{UV}^2} \tag{41}$$

As $r_0 \lesssim r_{UV}$ and $L > l_s$, in the end warping and wrapped volumes conspire so that only the orbifold and wrapping numbers can help to increase f. With large wrapping numbers and orbifold numbers, moderate[9] super-Planckian decay constants are attainable, arguably at the limits of the supergravity approximation and consistently with the Weak Gravity Conjecture. For example, taking $g_s \sim 0.3$, $L \sim \sqrt{3}\, l_s$, $r_{UV} \sim r_0 \sim 30\, l_s$, $n \sim 30$ and $k \sim 100$, we obtain $f \sim 3.4\, M_{Pl}$, $M_s \sim 3.1 \times 10^{-3} M_{Pl}$ and $\Lambda = 7.1 \times 10^{-2} M_{Pl}$. Beware again, however, that the backreaction of the D5-brane may actually be larger than the estimate given in (36), as the brane is at the top of the throat, far from the source N D3-branes. Also, the string scale against results to be quite low (cf. eq. (1)).

More complicated warped geometries may allow for more reliable and even larger axion decay constants. For example, a D5-brane moving in the warped resolved conifold was used in [40] to obtain a super-Planckian decay constant, and proposed as a model for Natural Inflation. The 10D metric for the warped resolved conifold is [55,56]:

$$ds_{10}^2 = h(r,\theta_2)^{-1/2} ds_4^2 + h(r,\theta_2)^{1/2} ds_6^2\,, \tag{42}$$

with

$$\begin{aligned} ds_6^2 = \kappa^{-1}(r)dr^2 &+ \frac{1}{6}r^2\left(d\theta_1^2 + \sin^2\theta_1 d\phi_1^2\right) \\ &+ \frac{1}{6}\left(r^2 + 6u^2\right)\left(d\theta_2^2 + \sin^2\theta_2 d\phi_2^2\right) \\ &+ \frac{1}{9}\kappa(r) r^2\left(d\psi^2 + \cos\theta_1 d\phi_1 + \cos\theta_2 d\phi_2\right)^2 , \end{aligned} \tag{43}$$

where

$$\kappa(r) = \frac{r^2 + 9u^2}{r^2 + 6u^2} \tag{44}$$

and u is the resolution parameter. The warp factor $h(r,\theta_2)$ is given by:

$$h(r,\theta_2) = L^4 \sum_{l=0}^{\infty}(2l+1) H_l(r) P_l(\cos\theta_2)\,, \tag{45}$$

[6] Note that for $p = 3$, the warp factor cancels out in f/M_{Pl}, eq. (28), and cannot be used to make f/M_{Pl} large.

[7] The volume orthogonal to the D5-brane embedded into spacetime with the static gauge is obtained by setting $\theta_1 = \phi_1 = 0$.

[8] Note that this condition is easier to achieve than the one required in [39], as in the latter the backreaction is enhanced by the brane's velocity.

[9] Larger values for f/M_{Pl} are possible compared to [39] by considering the brane moving in an angular direction rather than the radial direction.

where

$$H_l(r) = \frac{2}{9u^2}\frac{C_\beta}{r^{2+2\beta}}\,{}_2F_1\left(\beta, 1+\beta, 1+2\beta; -\frac{9u^2}{r^2}\right) \tag{46}$$

and

$$C_\beta = \frac{(3u)^{2\beta}\Gamma(1+\beta)^2}{\Gamma(1+2\beta)} \quad \text{and} \quad \beta = \sqrt{1+\frac{3}{2}l(l+1)}\,. \tag{47}$$

The warped geometry is sourced by N D3-branes placed at the tip of the throat, at $r = 0$ and $\theta_2 = 0$. Very close to the branes, at $\theta_2 = 0$ and small r, the geometry corresponds to $adS_5 \times S^5$ with $h = L^4/r^4$ and L given in (34). Further away from the branes, but still at small $r < u$, the warp factor is $h = L^4/(ru)^2$. Taking the smeared limit for the N source D3-branes, the warp factor can be approximated by the $l = 0$ mode as [55,56]:

$$h(r) = \frac{2L^4}{9r^2u^2} - \frac{2L^4}{81u^4}\log\left(1+\frac{9u^2}{r^2}\right). \tag{48}$$

Using this approximation to compute the warped volume, one obtains

$$V_6^w = s\frac{64\pi^3}{81k}L^4 r_{UV}^2 \tag{49}$$

with

$$s = 2 + \frac{1}{9}\frac{r_{UV}^2}{u^2}\left(1-\log\left(1+\frac{9u^2}{r_{UV}^2}\right)\right) - \frac{1}{81}\frac{r_{UV}^4}{u^4}\log\left(1+\frac{9u^2}{r_{UV}^2}\right). \tag{50}$$

The string scale is then given by:

$$M_s^2 = M_{Pl}^2\frac{81k}{4(2\pi)^4 s}\frac{g_s^2 l_s^6}{L^4 r_{UV}^2}. \tag{51}$$

In the near-tip region where the warp factor is $h_0 = L^4/(r_0u)^2$, $g_{\theta\theta} = u^2$ and the wrapped volume is $\mathcal{V}_2 l_s^2 = 4\pi u^2$, taking $u \sim r_{UV}$, the axion decay constant becomes[10]:

$$\frac{f^2}{M_{Pl}^2} = \frac{6561\,nk\,g_s}{32\pi^2(171-10\log 10)}\frac{l_s^2 u}{L^2 r_0}. \tag{52}$$

The possibility now emerges that small $r_0 < u$ could correspond to super-Planckian values [40].

The condition that the backreaction of the probe D5-brane be much smaller than that of the N D3-branes (26) can now be written as:

$$\frac{1}{2}\frac{n}{N}\frac{L^2}{l_s^2}\frac{r}{u}\sin\theta_1 \ll 1\,. \tag{53}$$

As the probe brane is close to the tip and the source N D3-branes, this should be a good indication of the reliability of the probe approximation. See also [40] for a more systematic check.

Also, the Weak Gravity Conjecture (21) now implies that:

$$1 < \frac{128(171-10\log 10)\pi^2}{6561}\frac{1}{g_s nk}\frac{r_0^2}{l_s^2}. \tag{54}$$

where we used that the masses of warped strings stretching between probe D-branes separated up the throat and around the throat are, respectively (cf. (37), (38)):

$$m = M_s^2\int_{r_0}^{r_{UV}} dr h^{1/4}\kappa^{-1/2}\,, \tag{55}$$

$$m = M_s^2\pi L\frac{u^{1/2}}{r_0^{1/2}}\,. \tag{56}$$

Notice again that the Weak Gravity Conjecture (54) ensures that the wrapping number, orbifold number and string couplings are not too large, as expected. It leads to a simple upper bound on the axion decay constant (see eq. (52)):

$$\frac{f^2}{M_{Pl}^2} < \frac{4r_0u}{L^2}. \tag{57}$$

It is now possible to achieve large f/M_{Pl} within the limits of perturbative control, small brane backreaction and consistently with the Weak Gravity Conjecture. For example, taking $g_s \sim 0.3$, $L \sim \sqrt{3}\,l_s$, $u \sim r_{UV} \sim 65\,l_s$, $r_0 \sim 2l_s$, $n \sim 10$ and $k \sim 10$, we find $f \sim 6.8\,M_{Pl}$, $M_s \sim 2.1\times 10^{-4}M_{Pl}$ and $\Lambda = 1.2\times 10^{-2}M_{Pl}$, consistently with (53) and (54).

Although we have seen that warping and non-trivial internal geometries allow one to consistently obtain axions with large axion decay constant in string theory, these axions cannot be used for large field inflation because the corresponding string scales are too small. Large field inflation for observable tensor modes implies that inflation occurs at least at GUT scale $M_{inf} \sim 10^{-2}M_{Pl}$ (see eq. (1)), so the string scale should be at least around $M_s \sim 10^{-1}M_{Pl}$, for a supergravity analysis to be valid during inflation.[11] In fact, as is well known, dimensional reduction gives the following relation between the string scale and Planck scale in any string model:

$$M_s^2 = M_{Pl}^2\frac{g_s^2}{4\pi\mathcal{V}_6^w}. \tag{58}$$

Thus a string scale $M_s \gtrsim 10^{-1}M_{Pl}$ is difficult to achieve within the limits of supergravity. This renders questionable all models of large field inflation in perturbative string theory, even those with sub-Planckian decay constants like axion monodromy.[12]

5. Discussion

We have shown that single open string axions coming from D-branes in warped geometries can enjoy super-Planckian axion decay constants consistently with the Weak Gravity Conjecture and within the limits of the supergravity approximation. The most explicit realisation is the model proposed by Kenton and Thomas [40] for Natural Inflation, using the position modulus of a wrapped probe D-brane moving in a compact angular direction near the tip

[10] This result differs from that in [40] since we use $h = \frac{L^4}{(ru)^2}$ valid for $\theta_2 \neq 0$, in contrast to $h = \frac{L^4}{r^4}$, used in [40]. This latter warp factor is also used to compute the warped volume of the throat up to the UV cutoff in [40].

[11] Strictly speaking, a hierarchy of scales $M_{inf} < M_c < M_s \lesssim M_{Pl}$, where M_c is the compactification scale, is needed in order to consistently work in the four dimensional, low energy supergravity limit of string theory [36,57], and moduli must also be stabilised.

[12] Factors like (2π) might help in general to achieve higher string scales within the perturbative limits, but explicit constructions enhance this problem. For example, in the original axion monodromy proposal [26], increasing $\mathcal{V}_6^w$ by using throats within throats to prevent brane anti-brane annihilation and suppress brane backreaction [58,59] will drive the string scale down. The Large Volume Scenario used in the D7-brane chaotic inflation model of [25] would also make a high string scale difficult to achieve. See also [60].

of the warped resolved conifold. Here, initial estimates of the backreaction of the brane, which is close to the source branes at the tip, also indicate that the probe approximation is reliable, though it would be important to verify this with explicit computation of the backreaction.

However, none of the constructions considered here can be used for large field inflation, as their string scales are too low compared to the inflationary scale (1). Indeed, we stress that one of the biggest challenges in realising large field inflationary models and observable gravitational waves within string theory is to identify appropriate compactifications with sufficiently high string scale, within the limits of four dimensional supergravity (or, alternatively, to evade the Lyth bound). Within the perturbative supergravity limit of string theory, dimensional reduction of the 10D Einstein–Hilbert term relates the string scale and Planck scale as in (58). Assuming $g_s \lesssim 0.3$ and $\mathcal{V}_6^w \gtrsim 1$ for the weak coupling and α' expansions, implies $M_s \lesssim 0.08 M_{Pl}$. Moreover, assuming a curvature scale l such that $(l_s/l)^2 \lesssim 0.3$ for a reliable α' expansion, implies $M_{kk} \lesssim 0.04 M_{Pl}$. Although numerical coefficients like factors of 2π could help, the required hierarchy $M_{inf} \ll M_{kk} \ll M_s$, seems very difficult to achieve for high scale inflation $M_{inf} \sim 0.01 M_{Pl}$. We could try to increase the string coupling and/or reduce the internal volumes to drive the string scale up, but then the relation (58) would not apply. In this case though, one could always dualize to an equivalent weak coupling, weak curvature description where (58) holds again, and return to the bounds $M_s \lesssim 0.08 M_{Pl}$, $M_{kk} \lesssim 0.04 M_{Pl}$.

Although the open string axions with large decay constant studied here cannot be used for large field inflation, they may give some insight into the Generalized Weak Gravity Conjecture. The D-brane axions have interpretation as (or are related to) gauge fields in higher dimensions, and so would be directly subject to the Weak Gravity Conjecture.[13] Furthermore, the Generalized Weak Gravity Conjecture, in its mild form, states that among any instantons, i, coupling to an axion with coupling f_i, there must exist at least one with action $S^i_{cl} \lesssim M_{Pl}/f_i$. The strong form states, additionally, that this instanton must be the one with smallest action.

However, we have seen that for open strings in warped geometries a large axion decay constant is generated within the limits of supergravity, where instanton effects must be exponentially suppressed. Indeed, the known instantons that couple to the Dp-brane Wilson line and position open string moduli are Euclidean Eq-branes wrapping $(q+1)$-cycles on the internal space. These instantons usually couple to the D-brane moduli with axion decay constant f, and their actions go as g_s^{-1} and the volume of the wrapped $(q+1)$-cycle.[14] Thus they have $S_{cl} > M_{Pl}/f$, and contributions are exponentially suppressed within the supergravity approximation. There then seems to be three likely possibilities: *(i)* In the presence of warping, the Eq-instantons couple to the open string axions with a suppressed effective axion coupling $f' < f$, such that they satisfy the Generalized Weak Gravity Conjecture with large action. *(ii)* There are new stringy instanton effects that couple to the D-brane moduli with a suppressed effective axion coupling $f' < f$, which satisfy the Generalized Weak Gravity Conjecture[15] with large action, $S'_{cl} \lesssim M_{Pl}/f'$, either mildly ($S'_{cl} > S_{cl}$) or strongly ($S'_{cl} < S_{cl}$). *(iii)* The Generalized Weak Gravity Conjecture is incorrect. We hope this letter will help to shed light on these issues.

[13] See the contemporaneous paper [61] for a further discussion on the Weak Gravity Conjecture and dimensional reduction and the Generalized Weak Gravity Conjecture for axions.

[14] See [62] for an explicit computation of the instanton generated scalar potential for open string moduli in toroidal orbifold compactifications with magnetized D-branes.

[15] This is similar – but slightly different – to the loophole discussed in [14] to achieve Natural Inflation consistently with the strong form of the Generalized Weak Gravity Conjecture.

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